Impact of a light strange-beauty squark on B_s mixing and direct search

Abdesslam Arhrib, Chun-Khiang Chua, and Wei-Shu Hou
Department of Physics, National Taiwan University, Taipei, Taiwan 10764, Republic of China
(Received 12 March 2001; published 12 December 2001)

If one has Abelian flavor symmetry, s_R-b_R mixing could be near maximal. This can drive a "strange-beauty" squark ($\widetilde{sb}_1$) to be rather light, but still evade the $b\to s\gamma$ constraint. Low-energy constraints imply that all other superpartners are at TeV scale, except for a possibly light neutralino, $\tilde{\chi}_1^0$. Whether light or heavy, the $\widetilde{sb}_1$ can impact on the B_s system: Δm_{B_s} and the indirect CP phase, even for $B_s\to\phi\gamma$. A direct search is similar to the usual $\tilde{b}\to b\tilde{\chi}_1^0$, but existing bounds are weakened by the $\widetilde{sb}_1\to s\tilde{\chi}_1^0$ possibility. All of these effects could be studied soon at the Fermilab of Tevatron.

DOI: 10.1103/PhysRevD.65.017701 PACS number(s): 12.60.Jv, 11.30.Er, 11.30.Hv, 13.25.Hw

The source of CP violation within the standard model (SM) rests in the flavor sector, which is not well understood. With three quark generations, we have six masses, three mixing angles, and a unique CP phase in the Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix V. The left-handed nature of weak dynamics screens out the mixings and CP phases (no longer unique) in the right-handed quark sector. The actual number of flavor parameters is much larger than meets the eye.

The observed quark masses and mixings do, however, exhibit an intriguing hierarchical pattern in powers of $\lambda \equiv |V_{us}|$, hinting at a possible underlying symmetry [1]. If this "horizontal" or flavor symmetry is Abelian, then s_R-b_R mixing would be near maximal [2,3], although still hidden from view. If supersymmetry (SUSY) is also realized, $\tilde{s}_R-\tilde{b}_R$ squark mixing could then be near maximal. This could generate observable effects in $b\to s$ transitions even if squark masses are at TeV scale [2,3]. Furthermore, one of the squarks, the "strange-beauty" squark $\widetilde{sb}_1$, could be driven by this large flavor violation to be considerably below the other squarks [3]. Whether we have a light $\widetilde{sb}_1$ squark or not, it is of great interest since the current bound on Δm_{B_s} [4] indicates that it could be larger than SM expectations.

In this Brief Report we point out that a light $\widetilde{sb}_1$ squark is allowed by the $b\to s\gamma$ constraint. We explore the implications of large $\tilde{s}_R$-$\tilde{b}_R$ mixing on B_s-$\overline{B_s}$ mixing and its CP phase Φ_{B_s}. In case of light $\widetilde{sb}_1$, we briefly comment on direct search. All these effects can be covered at the Tevatron Run II, which has just started. Mixing dependent CP violation in $B_s\to\phi\gamma$ decay can also be studied in the future. We stress that, besides the assumptions of Abelian flavor symmetry and SUSY, the quark mixing and CP phase we study are on similar footing as the usual CKM matrix.

Horizontal models try to explain the mass and mixing hierarchies by powers of $\lambda\sim\langle S\rangle/M$, where $\langle S\rangle$ is the expectation of a scalar field S and M is a high scale. For Abelian symmetries, the commuting nature of horizontal charges in general gives $M_{ij}M_{ji}\sim M_{ii}M_{jj}$ (i, j not summed), where "$\sim$" indicates approximate. This allows one to determine, e.g., M_d^{32} from our knowledge of $M_d^{22}\sim m_s\sim\lambda^2 m_b$, $M_d^{23}\sim V_{23}m_b\sim\lambda^2 m_b$ and $M_d^{33}\sim m_b$. Hence [2]

$$\hat{M}_d=\frac{M_d}{m_b}\sim\begin{bmatrix}\lambda^4 & [\lambda^3] & [\lambda^3]\\ [\lambda^3] & \lambda^2 & \lambda^2\\ [\lambda] & 1 & 1\end{bmatrix}, \tag{1}$$

and similarly for M_u; the $[\cdots]$ terms would be set to zero as explained shortly. Diagonalizing M_d by a biunitary D_L and D_R transform, $D_R^{23}\sim 1$ is clearly the largest mixing element, but its effect is hidden within SM.

Taking SUSY as commuting with the horizontal symmetry, the squark mass matrices are fixed by the common horizontal charge of the chiral supermultiplet. We take the usual approach that squarks are almost degenerate with common scale $\tilde{m}$. From Eq. (1) one finds that $(\tilde{M}_d^2)_{LR}=(\tilde{M}_d^2)_{RL}^\dagger\sim\tilde{m}M_d$, $(\tilde{M}_d^2)_{LL}\sim\tilde{m}^2V$, while

$$(\tilde{M}_d^2)_{RR}\sim\tilde{m}^2\begin{bmatrix}1 & [\lambda] & [\lambda]\\ [\lambda] & 1 & 1\\ [\lambda] & 1 & 1\end{bmatrix}, \tag{2}$$

where $(\tilde{M}_d^2)_{RR}^{23,32}\sim\tilde{m}^2$ if s_R and b_R have the same horizontal charge(s), hence comparable to $(\tilde{M}_d^2)_{RR}^{22,33}\sim\tilde{m}^2$.

We are interested in the impact of $(\tilde{M}_d^2)_{RR}^{23,32}$. It is known that four texture zeros are needed [3] to fully evade the Δm_K and ε_K constraints. Hence, we choose horizontal charges such that the $[\cdots]$ terms in Eqs. (1) and (2) are all set to zero, which is achievable under a U(1)×U(1) or higher horizontal group. With the d quark thus decoupled, one is safe from all known low-energy constraints. However, one needs $(\tilde{M}_u^2)_{LL}^{12}\sim\lambda\tilde{m}^2$ to account for V_{us} [2]. It is intriguing that $\tilde{m},m_{\tilde{g}}\sim$TeV brings [3] Δm_D right into the ballpark of current [5] experimental sensitivities. This sets the scale for $\tilde{m}$ and $m_{\tilde{g}}$, for if they were lighter, Δm_D would be too large. Similarly, Δm_K constrains $\tilde{u}_L$, $\tilde{c}_L$, and $\tilde{\chi}^\pm$ loops, implying also [3] that squarks are at TeV scale, while the W-ino part of the chargino is heavier than 500 GeV.

With d-flavor decoupled, the s-b part of $\tilde{M}_{RR}^2$ in Eq. (2) appears "democratic." More explicitly, one has

$$\tilde{M}_{RR}^{2(sb)}=\begin{bmatrix}\tilde{m}_{22}^2 & \tilde{m}_{23}^2e^{-i\sigma}\\ \tilde{m}_{23}^2e^{i\sigma} & \tilde{m}_{33}^2\end{bmatrix}=R\begin{bmatrix}\tilde{m}_1^2 & 0\\ 0 & \tilde{m}_2^2\end{bmatrix}R^\dagger, \tag{3}$$

in quark mass basis, where $\tilde{m}_{ij}^2\sim\tilde{m}^2$ are all >0, and

$$R=\begin{bmatrix}c_\theta & s_\theta\\ -s_\theta e^{i\sigma} & c_\theta e^{i\sigma}\end{bmatrix}. \tag{4}$$

The phase in R absorbs the σ phase in $\tilde{M}_{RR}^{2(sb)}$. By nature of $\tilde{m}_{23}^2 \sim \tilde{m}_{22,33}^2$, in general we have one suppressed eigenvalue $\tilde{m}_1^2$ due to level splitting, where θ is a measure of the relative weight of $\tilde{m}_{23}^2$ vs $\tilde{m}_{33}^2 - \tilde{m}_{22}^2$. Since our case corresponds to $\tilde{m}_{22}^2 \simeq \tilde{m}_{33}^2 \simeq \tilde{m}_{23}^2 \simeq \tilde{m}^2$ because of Eqs. (1) and (2), near maximal mixing is implied. The eigenstates hence carry both s and b flavors and are called the strange-beauty squarks $\widetilde{sb}_{1,2}$. Without much loss of generality, we take $\tilde{m}_{22}^2 = \tilde{m}_{33}^2 = \tilde{m}^2$ (so $\theta = \pi/4$ and $\tilde{m}_1^2 + \tilde{m}_2^2 = 2\tilde{m}^2$) and consider the ratio $\tilde{m}_{23}^2/\tilde{m}^2 \equiv 1 - \delta \simeq 1$. The squark mass eigenvalues must be positive to preserve color symmetry, hence $\delta > 0$ is required. For small δ, we have $\tilde{m}_1^2 \cong \delta \tilde{m}^2$ and $\tilde{m}_2^2 \cong (2-\delta)\tilde{m}^2$. Thus, with some tuning, $\widetilde{sb}_1$ can become quite light, i.e., $\tilde{m}_1^2 \ll \tilde{m}_2^2 \simeq 2\tilde{m}^2$, the driving force being the large $(\tilde{M}_d^2)_{RR}^{23,32}/\tilde{m}^2 \sim 1$ in Eqs. (2) and (3). We note that, assuming $\tilde{m} \sim 2$ TeV, tuning δ to λ^2, λ^3, and λ^4 give $\tilde{m}_1 = 440$, 206, and 97 GeV; for $\tilde{m} \sim 1$ TeV, $\delta = \lambda$, λ^2, and λ^3 gives $\tilde{m}_1 = 470$,220, and 103 GeV, respectively. In the following, we limit ourselves to $\tilde{m}_1 \geqslant 100$ GeV.

In addition to concerns about tuning, the pressing question is that a light $\widetilde{sb}_1$ driven by large strange-beauty mixing seems particularly dangerous in face of the $b \to s\gamma$ constraint. As shown in [3], heavy squark and gluino loops are suppressed by $1/G_F\tilde{m}^2$ compared to the SM contribution, such that $b \to s\gamma$ rate is hardly affected. It is interesting that, *even with $\widetilde{sb}_1$ as light as 100 GeV, the $b \to s\gamma$ constraint is still rather accommodating.*

The calculation of short distance coefficients is done following [6]. The expressions for Wilson coefficients together with their renormalization-group equations (RGE) can be found in [7,8]. Our model gives large RR and RL mixings, while LL and LR mixings are suppressed by λ^2. In terms of the loop-induced effective $bs\gamma$ couplings $m_b \bar{s}[C_7 R + C_7' L]\sigma_{\mu\nu} F^{\mu\nu} b$, it is C_7' that receives larger contributions. This in itself provides some protection, since C_7' is not generated in SM ($C_7^{\rm SM} \simeq -0.31$); hence our SUSY effects enter the $b \to s\gamma$ rate only quadratically.

We find that, although RL mixing is suppressed by $m_b/\tilde{m}$, its effect dominates over the RR contribution for $\cos\sigma < 0$. Let us first show that C_{7RR}' is finite and suppressed by $m_{\tilde{g}}^2$ in the $\tilde{m}_1^2 \to 0$ limit. By direct computation, one finds that the $\widetilde{sb}$-$\tilde{g}$ loop contribution to $m_b C_{7RR}'$ is proportional to

$$\int dk^2 \frac{k^4 m_b \tilde{m}_1^2 c_\theta s_\theta e^{-i\sigma}}{(k^2+m_{\tilde{g}}^2)(k^2+\tilde{m}_1^2)^4} - (\tilde{m}_1^2 \to \tilde{m}_2^2), \qquad (5)$$

where "super-Glashow-Iliopoulos-Maieesci (GIM)" cancellation is ensured and the $\widetilde{sb}_2$ term decouples for heavy $\tilde{m}_2^2$. Since RR mixing is chiral conserving, a factor of m_b is needed, while $\tilde{m}_1^2 c_\theta s_\theta e^{-i\sigma}$ is from $(\tilde{M}_d^2)_{RR}^{23}$. The integral is clearly finite in the $\tilde{m}_1^2 \to 0$ limit. Using formulas from [8], we find $C_{7RR}'(M_{\rm SUSY}) \simeq -0.1\, c_\theta s_\theta e^{-i\sigma} (0.8\ {\rm TeV}/m_{\tilde{g}})^2$ for the maximal super-GIM breaking (small $\tilde{m}_1^2$, large $\tilde{m}_2^2$) case. The shift in the $b \to s\gamma$ rate is $\lesssim 2\%$ for $\tilde{m} > m_{\tilde{g}} > 0.8$ TeV.

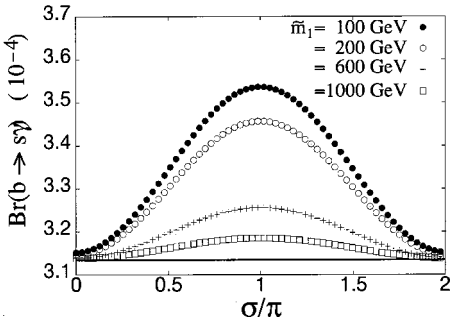

FIG. 1. $b \to s\gamma$ vs CP phase σ including both SM and SUSY effects, for $m_{\tilde{g}}$, $\tilde{m} = 0.8$, 2 TeV and several strange-beauty squark mass ($\tilde{m}_1 \equiv m_{\widetilde{sb}_1}$) values. The horizontal line indicates the SM expectation.

The $m_b/\tilde{m}$ suppression of the RL contribution is compensated by a chiral enhancement factor [9] $m_{\tilde{g}}/m_b$. With $\tilde{b}_L \to \widetilde{sb}_{1R}$ mixing and $\tilde{b}_L$ heavy, the $s_\theta e^{-i\sigma}$ factor in Eq. (5) is replaced by $(R^\dagger \tilde{M}_{RL}^2)^{23}/\tilde{m}^2 \sim (c_\theta - s_\theta e^{-i\sigma}) m_b/\tilde{m}$, where R is given in Eq. (4), and we take $(\tilde{M}_d^2)_{RL}^{23,33} \sim m_b \tilde{m}$ as real. The factor $k^4/(k^2+\tilde{m}_1^2)^4$ in Eq. (5) is replaced by $k^2/(k^2+\tilde{m}_1^2)^3$ and the integral is still finite for $\tilde{m}_1^2 \to 0$. We find $C_{7RL}'(M_{\rm SUSY}) \simeq 0.12 c_\theta (c_\theta - s_\theta e^{-i\sigma})(1.6\ {\rm TeV}^2/\tilde{m} m_{\tilde{g}})$ for small $\tilde{m}_1^2$. Taking $s_\theta \sim 1/\sqrt{2}$, C_{7RL}' is subdominant for $\sigma \sim 0$, but dominates over C_{7RR}' for $\sigma \sim \pi$.

We illustrate in Fig. 1 the full gluino and neutralino loop effect on $b \to s\gamma$ rate vs CP phase σ, for $m_{\tilde{g}} = 0.8$ TeV and $\tilde{m} = 2$ TeV, with simplifying assumptions as stated above. Even for $\tilde{m}_1$ as light as 100 GeV, $b \to s\gamma$ is still [10] well within the allowed experimental range of $(3.15 \pm 0.54) \times 10^{-4}$ [11]. For heavy $\widetilde{sb}_1 \sim 1$ TeV, its effect becomes negligible, and the $b \to s\gamma$ rate approaches the SM value, as indicated by the horizontal line at $b \to s\gamma \sim 3.14 \times 10^{-4}$ for our parameter choice. The σ dependence can be understood through our earlier discussion of C_{7RR}' and C_{7RL}'. One can also easily check from the strength of $|C_7'|^2$ as seen in Fig. 1, that the LR mixing contribution $\delta C_{7\rm LR}$ is indeed subdominant even though it interferes with $C_7^{\rm SM}$, which is large.

It is intriguing that, although C_7' is subdominant compared to $C_7^{\rm SM}$, its strength is actually not small. That is, $|C_7'/C_7| \simeq 0.35 - 0.12$; hence $\sin 2\vartheta = 2|C_7 C_7'|/(|C_7|^2 + |C_7'|^2) \simeq 63\% - 22\%$ for $\tilde{m}_1 = 100 - 1000$ GeV. New physics effects [3,8] such as mixing dependent CP violation in $B^0 \to K_1^0(1270)\gamma$ could be of this order (though direct CP is small because $\delta C_{7\rm LR}$ is small), but detectability may be better in $B_s \to \phi\gamma$. "Wrong" Λ polarization in $\Lambda_b \to \Lambda\gamma$ could also be promising [8].

It is known that charged Higgs effects on $b \to s\gamma$ add constructively to the SM for all $\tan\beta$ [12], giving rise to a very stringent constraint on m_{H^+}. Our light $\widetilde{sb}_1$ only slightly worsens the situation.

Turning to charginos, as stated, the Δm_K constraint de-

mands that the *W*-ino part of chargino mass, controlled by M_2, should be larger than 500 GeV. We do not entertain a light top squark since we tacitly assume that flavor and SUSY scales are not too far apart [2], so the up squark mass average $\tilde{m}_u$ is also at $\tilde{m}\sim$ TeV. Thus, the charginos and the *W*-ino or Higgsinolike neutralinos are all at the TeV scale. This still leaves open the possibility of a light bino with mass controlled by M_1, which we call $\tilde{\chi}_1^0$. Interestingly, $b\to s\gamma$ is not very constraining here; we have taken the rather low mass value of $m_{\tilde{\chi}_1^0}=90$ GeV in Fig. 1, and find that its effect is still much smaller than the dominant gluino contribution. This is simply because of the much weaker *B*-ino coupling (hypercharge) to down sector compared with the strong gluino couplings.

Without necessarily advocating a light *B*-ino, we thus have a scenario where SUSY particles and exotic Higgs bosons are at TeV scale, except for a possibly *light neutralino* $\tilde{\chi}_1^0$ *that is largely B-ino,* and a *light strange-beauty squark* $\widetilde{sb}_1$ *with mass driven low by flavor violation!*

One may worry that large $\tilde{q}$-$\widetilde{sb}_1$(or$\tilde{\chi}^-$-$\tilde{\chi}_1^0$) splittings may violate $\delta\rho$ constraint. We first note that $\delta\rho$ picks up corrections to isovector gauge boson self-energy diagrams. Because the isovector gauge interaction is left handed, contributions from right-handed squarks are transmitted through LR mixing [13]. However, this is suppressed in our case by $\tilde{M}_{LR}^2/\tilde{m}^2\sim m_{s,b}/\tilde{m}\sim 10^{-4}-10^{-3}$ [14]. $\delta\rho$ can constrain only mass splittings in $\tilde{q}_L$, which are TeV scale particles and do not have large splittings, and thus the seemingly dangerous large splitting involving $\widetilde{sb}_1$ is safe from $\delta\rho$ constraint. We note in passing that our light $\widetilde{sb}_1$ can evade the R_b constraint also. The $\tilde{\chi}^0$-$\tilde{d}_j$ contribution to R_b is negligible [15] while $\tilde{\chi}^-$-$\tilde{t}$ gives sizable contribution only for light stop and light chargino, which is not the case in our model.

Large $\tilde{s}_R$-$\tilde{b}_R$ mixing, however, can easily impact on B_s-$\overline{B}_s$ mixing and its *CP* phase Φ_{B_s}, will soon be accessible at the Tevatron. Recall that $(\tilde{M}_d^2)_{RR}^{23}/(\tilde{M}_d^2)_{RR}^{13}\sim 1/\lambda\sim|V_{ts}/V_{td}|$ in Eq. (2), before setting [. . .] terms to zero. By simply scaling up the B_d mixing results of [3] for the $\tilde{d}_R$-$\tilde{b}_R$ mixing case, one sees that even for $\tilde{m}_1\sim$ TeV, its contribution to B_s mixing could be of the same order as SM. The dominant $\tilde{q}$-$\tilde{g}$ box diagrams involve two $\widetilde{sb}_1$, or one $\widetilde{sb}_1$ and one $\tilde{s}_L/\tilde{b}_L$ with $\tilde{s}_L$-$\tilde{b}_L$ mixing. The former generates effective coupling $\propto\tilde{C}_1\overline{s_R^\alpha}\gamma_\mu b_R^\alpha\overline{s_R^\beta}\gamma^\mu b_R^\beta$, while the latter $\propto C_{4(5)}\overline{s_R^\alpha}b_L^{\alpha(\beta)}\overline{s_L^\beta}b_R^{\beta(\alpha)}$, where $\tilde{C}_1\propto c_\theta^2 s_\theta^2 e^{-2i\sigma}, C_{4(5)}\propto\lambda^2 c_\theta s_\theta e^{-i\sigma}$ are known [6,16] functions of $m_{\tilde{g}}^2/\tilde{m}_1^2$ (simpler mass insertion formulas given in [3]). Because of a larger loop factor, the CKM suppressed $C_{4(5)}$ is comparable to $\tilde{C}_1$. Thus, the explicit σ-phase dependence of the mixing amplitude is (a, b, and c are real)

$$M_{12}\equiv|M_{12}|e^{2i\Phi_{B_s}}\cong ae^{-2i\sigma}+be^{-i\sigma}+c, \qquad (6)$$

where b (from $C_{4(5)}$) and c (from SM) differ in sign.

Using the RGE evolution from [17] and $f_{B_s}^2 B_{B_s}=(240\ \text{MeV})^2$, we find $\Delta m_{B_s}^{\rm SM}\simeq 14.9\ \text{ps}^{-1}$ with vanishing $\sin 2\Phi_{B_s}^{\rm SM}$. For illustration, in Fig. 2 we plot Δm_{B_s} and $\sin 2\Phi_{B_s}$ vs σ for $\tilde{m}_1=1.2$ TeV, average squark mass $\tilde{m}=2$ TeV and $m_{\tilde{g}}=1$, 2, and 3 TeV. Even for heavy $\widetilde{sb}_1$ at TeV scale, the SUSY contribution can be comparable to the SM effect. For $m_{\tilde{g}}=1$ TeV$<\tilde{m}_1$, Δm_{B_s} can reach twice the SM value around $\sigma\sim\pi$. For heavier $m_{\tilde{g}}$, Δm_{B_s} can reach only 22 (18) ps^{-1} for $m_{\tilde{g}}=2(3)$ TeV. Destructive interference between SM and SUSY for $\cos\sigma>0$ (where $\cos 2\sigma$ modulation can be seen) would give $\Delta m_{B_s}<\Delta m_{B_s}^{\rm SM}$ hence disfavored. Thus, for the $\widetilde{sb}_1\sim$ TeV scenario, $\cos\sigma<0$ is preferred. Similarly, $|\sin 2\Phi_{B_s}|$ can reach 50%$-$75%, vanishes at $\sigma=\pi$, and has a smaller range for heavier $m_{\tilde{g}}$. If Δm_{B_s} is only slightly above SM expectation, it could be uncovered at the Tevatron in a couple of years. One could then find $\sin 2\Phi_{B_s}\neq 0$ and indirect *CP* in $B_s\to\phi\gamma$, but no sign of SUSY particles since the scale is at TeV.

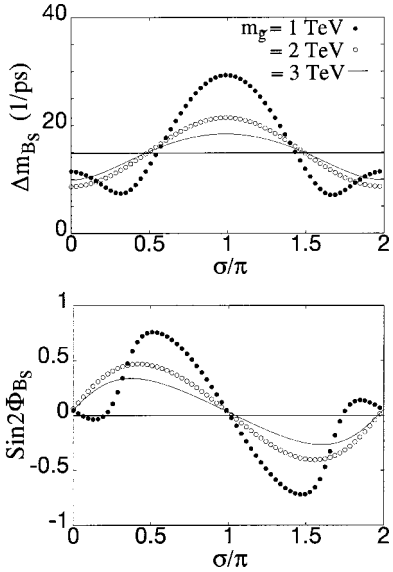

FIG. 2. Δm_{B_s} and $\sin 2\Phi_{B_s}$ vs σ for $\tilde{m}_1,\tilde{m}=1.2$, 2 TeV, and $m_{\tilde{g}}=1$, 2, 3 TeV. The horizontal line is SM expectation.

The light $\widetilde{sb}_1$ case allows greater range. We plot Δm_{B_s} and $\sin 2\Phi_{B_s}$ vs σ in Fig. 3, for $m_{\tilde{g}},\tilde{m}=0.8$, 2 TeV, and $\tilde{m}_1=100$, 200 and 600 GeV. The $\tilde{m}_1=600$ GeV case is similar to Fig. 2, except that $a+b$ in Eq. (6) is of the same sign as c. For lower $\tilde{m}_1$, the strength of b increases monotonically and is stronger than c, while a first drops slowly, resulting in an accidental cancellation of Δm_{B_s} at $\sigma=0$ for $\tilde{m}_1\sim 200$ GeV. Below this, a flips sign and changes rapidly, and together with b they overwhelm c. Thus, for $\tilde{m}_1\lesssim 130$ GeV, one develops a dip rather than maximum at $\sigma\sim\pi$, as shown for the $\tilde{m}_1=100$ GeV case.

It is interesting that Δm_{B_s} hovers not far above 15 ps^{-1} for both a broad range of $\tilde{m}_1\gtrsim 250$ GeV and $\cos\sigma>0$, and the intriguing case of a rather light (<100 GeV)$\widetilde{sb}_1$ for phase $\sigma\sim\pi$. For such Δm_{B_s} values, measurement would be swift, with good prospects for $\sin 2\Phi_{B_s}$, which clearly covers

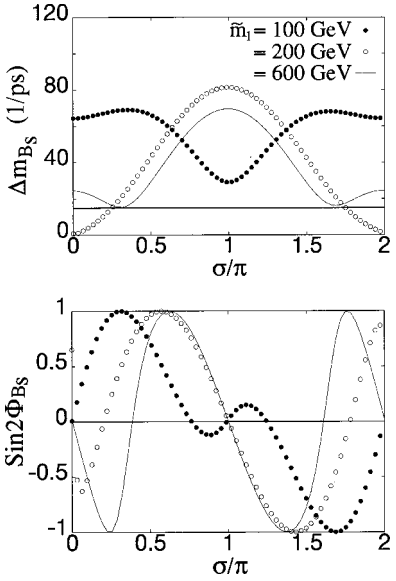

FIG. 3. Δm_{B_s} and $\sin 2\Phi_{B_s}$ vs σ for $m_{\tilde{g}}, \tilde{m} = 0.8$, 2 TeV, and three $\tilde{m}_1$ values. The horizontal line is SM expectation.

the full range between ± 1, with a $\sin 2\sigma$ modulation over the basic $\sin\sigma$ dependence. However, Δm_{B_s} can also easily reach beyond 40 ps^{-1}, whether $\widetilde{sb}_1$ is heavy or light, and measurement would then take a while. This in itself would indicate new physics, but $\sin 2\Phi_{B_s}$ measurement becomes difficult. For confirming evidence, one would have to search for C_7' effects in $b \to s\gamma$, such as indirect CP in $B_d \to K_1^0 \gamma$ or "wrong" Λ polarization in $\Lambda_b \to \Lambda\gamma$.

Whether Δm_{B_s} (and $\sin 2\Phi_{B_s}$) is measured soon or not, it is imperative to check whether there is a $\widetilde{sb}_1$ squark below a couple hundred GeV. In the usual SUSY scenario, because of heaviness of the top quark, one could have a light stop by RGE evolution from very high scale, or by having large (flavor blind) LR mixing. One could also have a light sbottom if $\tan\beta$ is large. This has motivated the experimental search [18] via $\tilde{b}_1 \to b\tilde{\chi}_1^0$ assuming that $\tilde{\chi}_1^0$, if not the lightest SUSY particle (LSP), is lighter than $\tilde{b}_1$. The signature is two b jets plus missing energy. In order to distinguish sbottom from stop, b tagging is necessary since loop-induced $\tilde{t}_1 \to c\tilde{\chi}_1^0$ leads to a similar signature. In our case, all squarks including stop are at TeV scale, except $\widetilde{sb}_1$ which becomes light because of large flavor violation, *without the need for large* $\tan\beta$. Since $\widetilde{sb}_1$ is a mixture of $\tilde{s}_R$ and $\tilde{b}_R$, both decays $\widetilde{sb}_1 \to b\tilde{\chi}_1^0$, $s\tilde{\chi}_1^0$ are important, and the b-tagging efficiency is diluted. Thus, the standard sbottom search bound would weaken. In any case, if a light sbottom is found, one would have to check against production cross section vs theory expectations from mass measurement, to determine whether it is the standard $\tilde{b}_1$ or the $\widetilde{sb}_1$. In case $\tilde{\chi}_1^0$ is heavier than $\widetilde{sb}_1$, the LSP would likely be some sneutrino, and the decay $\widetilde{sb}_1 \to b\tilde{\nu}\nu, s\tilde{\nu}\nu$ via virtual $\tilde{\chi}_1^0$ (hypercharge coupling) has a similar signature.

In conclusion, flavor violation in the $\tilde{s}_R$-$\tilde{b}_R$ squark sector could be uniquely large if one has an underlying *Abelian flavor symmetry*, which are both inspired by the hierarchical patterns of quark masses and mixings. With SUSY above TeV scale, this could evade low-energy constraints, including $b \to s\gamma$, but modify B_s mixing and generate $\sin 2\Phi_{B_s} \neq 0$. It is intriguing that the strange-beauty squark $\widetilde{sb}_1$ *could be driven light by the large flavor violation itself*. Both a light $\widetilde{sb}_1$ *and* a light B-ino-like neutralino $\tilde{\chi}_1^0$ can survive the $b \to s\gamma$ constraint. This would not only further enrich B_s physics, but can also be directly probed via $\widetilde{sb}_1 \to b\tilde{\chi}_1^0, s\tilde{\chi}_1^0$, which extends the standard $\tilde{b} \to b\tilde{\chi}_1^0$ search scenario.

This work is supported in part by NSC-89-2112-M-002-063, NSC-89-2811-M-002-0086, and 0129, the MOE CosPA Project, and the BCP Topical Program of NCTS.

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